

The η NN Coupling in Eta Photoproduction

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Abstract

Low-energy eta photoproduction on the nucleon is studied in an effective Lagrangian approach that contains Born terms, vector meson and nucleon resonance contributions. The resonance sector includes the $S_{11}(1535)$, $P_{11}(1440)$ and $D_{13}(1520)$ states whose couplings are fixed by independent electromagnetic and hadronic data. The available (γ, η) data are employed to discuss the difference between pseudoscalar and pseudovector Born terms and to determine the magnitude of the ηNN coupling constant. We present multipoles that are most sensitive to the various model ingredients and demonstrate how these multipoles may be accessed in polarization observables. Cross section

calculations are presented for eta photoproduction on light nuclei.

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I. INTRODUCTION

Over the last several years there has been renewed interest in the production of η -mesons with protons, pions and electrons and their interaction with nucleons and nuclei. One of the first η -nuclear experiments, performed at SATURNE in 1988 [1], reported surprisingly large eta production rates near threshold in the reaction $d(p,\eta)^3\text{He}$. These large cross sections permitted not only a more precise determination of the η -mass [2] but were also used to perform rare decay measurements of the eta [3]. Additional experiments involving pion induced eta production were performed at Los Alamos [4]. Again, the experimental cross sections at threshold region of the reaction $^3\text{He}(\pi^-, \eta)^3\text{H}$ are above the theoretical calculations [4,5].

The advent of high duty-cycle electron accelerators opens for the first time the opportunity to study the reactions $N(\gamma, \eta)N$ and $N(e, e'\eta)N$ in greater detail. Our present knowledge of the (γ, η) process is based solely on some old measurements around 20 years ago [6], along with very few more recent data from Bates [7] and Tokyo [8]. With the recent completion of the electron accelerators at Mainz (MAMI B) and Bonn (ELSA) and the construction of new spectrometers and detectors it is now possible to measure eta photoproduction from threshold at 707 MeV up to 850 MeV at Mainz and even higher energies at Bonn with a precision similar to the one obtained in pion photoproduction experiments. A large amount of data has already been taken and is currently being analyzed [9,10].

Unlike pion photoproduction, low energy theorems (*LET*) cannot be derived for eta photoproduction for the following three reasons: (*i*) The expansion parameter $\mu = m_\eta/m_N \approx 0.6$ is too large to provide convergence up to order μ^2 ; (*ii*) due to large $\eta - \eta'$ mixing with a mixing angle of about 20° and a non-conserved axial singlet current A_0^μ for the η' , there is no *PCAC* theorem for eta mesons; (*iii*) there are nucleon resonances, mainly the $S_{11}(1535)$ close at threshold ($W_{thr.} = 1486$ MeV) strongly violating the condition that the internal excitation energy must be larger than the mass of the meson [11,12].

Nucleon resonance excitation is the dominant reaction process in (γ, η) . In contrast to pions which will excite $\Delta(T = 3/2)$ as well as $N^*(T = 1/2)$ resonances, the η meson will

only appear in the decay of N^* resonances with $T = 1/2$. In the low-energy region this is dominantly the $S_{11}(1535)$ state that decays in 45-55% into ηN , the only nucleon resonance with such a strong branching ratio in the η channel. This result is even more surprising as a near-by resonance of similar structure, the $S_{11}(1650)$ has a branching ratio of only 1.5%. This " η puzzle" is not yet understood in quark models of the nucleon.

Most attempts to describe eta photoproduction on the nucleon [13,14] have involved Breit-Wigner functions for the resonances and either phenomenology or a Lagrangian approach to model the background. These models which contain a large number of free parameters were then adjusted to reproduce the few available data. In a very different approach, Ref. [15] derived a dynamical model which employs $\pi N \rightarrow \pi N$, $\pi N \rightarrow \pi\pi N$ and $\pi^- p \rightarrow \eta n$ to fix the hadronic vertex as well as the propagators and the $\gamma N \rightarrow \pi N$ to construct the electromagnetic vertex. This calculation represents a prediction rather than a fit to the $\gamma N \rightarrow \eta N$ reaction.

Here we extend the model of Ref. [15] by taking into account the background from s, u -channel nucleon Born terms and ρ, ω exchange in the t -channel. Since the resonance sector is fixed in our approach and also the vector meson couplings can be obtained from independent sources, we can use this model to extract information on the ηNN coupling. Furthermore, we apply the operator to elastic eta photoproduction on the very light nuclei, d , ^3He , ^3H and ^4He . Due to spin and isospin selection rules, measuring a combination of processes on these light nuclei with well-known nuclear structure should hopefully allow us a complete determination of the individual (γ, η) multipoles for protons and neutrons.

In Sec. II we shortly summarize the resonance model of Ref. [15] and describe our full (γ, η) operator. Sec. III contains our discussion of the elementary ηNN vertex. The possibility to get new information about nucleon resonances in the eta channel with the help of polarization observables is discussed in Sec. IV. The formalism for eta photoproduction on nuclei is derived in a coupled channel framework in Sec. V and we present predictions for differential cross sections of elastic (γ, η) on d , ^3He , ^3H and ^4He . In Sec. VI we summarize our findings and present a brief outlook. In the Appendix we present the definitions of the

16 polarization observables for photoproduction of pseudoscalar mesons and give expansions in CGLN amplitudes and dominant multipoles.

II. ETA PHOTOPRODUCTION ON THE NUCLEON

The dynamical model of Bennhold and Tanabe [15] is based on the observation that near the η production threshold three nucleon resonances $P_{11}(1440)$, $D_{13}(1520)$ and $S_{11}(1535)$ play an important role. Assuming an isobar model for each partial wave the transition amplitude can be written as

$$t_{ij}(W) = f_i^\dagger D^{-1}(W) f_j, \quad (1)$$

where W is the invariant energy and $i, j = \pi, \eta$ denotes the πN and ηN channels, respectively. The vertex functions f_i are parametrized with coupling strengths and formfactors and the N^* propagators are given by

$$D(W) = W - m_0 - \Sigma_\pi(W) - \Sigma_\eta(W) + \frac{i}{2}\Gamma_{\pi\pi}(W) \quad (2)$$

with the bare resonance mass m_0 .

The self-energy Σ associated with the πN and ηN intermediate states is given by

$$\Sigma_i(W) = \int_0^\infty \frac{q^2 dq}{(2\pi)^3} \frac{M}{2w_i(q)E_N(q)} \left(\frac{q}{m_i}\right)^{2l} \frac{g_i^2(1+q^2/\Lambda_i^2)^{-2-l}}{W - w_i(q) - E_N(q) + i\epsilon} \quad (3)$$

with $w_i(q) = \sqrt{m_i^2 + q^2}$, $E_N(q) = \sqrt{M^2 + q^2}$ and M denoting the nucleon mass. The two-pion decay width $\Gamma_{\pi\pi}$ is parametrized with one free parameter. The six parameters in this approach have been determined for each partial wave by a least-squares fit to all data of the reactions $\pi N \rightarrow \pi N$, $\pi N \rightarrow \pi\pi N$ and $\pi^- p \rightarrow \eta n$ and can be found in Ref. [15].

In order to use this operator more conveniently, especially in nuclear applications with multidimensional integrals, we have obtained simple parametrizations of the self-energy Σ , Eq. (3), in very good agreement with the exact numerical values.

$$Re \Sigma = a + (b_1\sqrt{x} + b_2x^2)\Theta(-x) + (c_1x + c_2x^2)\Theta(x), \quad (4a)$$

$$Im \Sigma = (d_1\sqrt{x} + d_2x + d_3x^2)\Theta(x) \quad (4b)$$

with $x = (W - M - m_i)/m_\pi$, $i = \pi, \eta$ and the step function $\Theta(x)$. The parameters are given in Table I. Finally the decay width in the 2π -channel is given by

$$\Gamma_{\pi\pi}(W) = \gamma x \Theta(x), \quad x = (W - M - 2m_\pi)/m_\pi \quad (5)$$

with $\gamma(S_{11}) = 4.3 \text{ MeV}$, $\gamma(P_{11}) = 80.3 \text{ MeV}$ and $\gamma(D_{13}) = 24.2 \text{ MeV}$.

With the hadronic vertices and propagators determined, the photoproduction amplitudes for (γ, π) and (γ, η) are given by

$$t_{i\gamma}(W) = V_{i\gamma}^B(W) + f_i^\dagger D^{-1}(W) \tilde{f}_\gamma, \quad (6)$$

where $V_{i\gamma}^B$ are the Born terms and \tilde{f}_γ is the electromagnetic vertex. The latter was determined by using pion photoproduction data. By this way no free parameters are introduced in the (γ, η) process. Since Ref. [15] neglected the Born terms in the η -channel, $V_{\eta\gamma}^B \equiv 0$, the model consisted of four (γ, η) multipoles only, the $S_{11}(1535)$ appears in the dominant E_{0+} , the $P_{11}(1440)$ in the M_{1-} and the $D_{13}(1520)$ in the E_{2-} and M_{2-} .

While neglecting the (γ, η) Born terms was within the uncertainties of the older experimental data for the proton, they play a more important role when comparison with better data becomes possible. Furthermore, including the background properly becomes necessary in nuclear reactions like the coherent η photoproduction on ^4He , where the dominant excitation of the S_{11} resonance is forbidden.

The evaluation of the background terms is straightforward and in complete analogy to (γ, π^0) except for the fact that the η is an isoscalar meson and two types of ηNN couplings are possible: Pseudovector (PV) and pseudoscalar (PS). The latter one is not ruled out by LET as in the case of (γ, π) .

The effective Lagrangians for both types of the ηNN coupling are given by

$$\mathcal{L}_{\eta NN}^{PS} = -ig_\eta \bar{\psi} \gamma_5 \psi \phi_\eta, \quad \mathcal{L}_{\eta NN}^{PV} = \frac{g_\eta}{2M} \bar{\psi} \gamma_\mu \gamma_5 \psi \partial^\mu \phi_\eta. \quad (7)$$

With the electromagnetic Lagrangian

$$\mathcal{L}_{\gamma NN}^{PS} = -e \bar{\psi} \gamma_\mu \frac{1 + \tau_0}{2} \psi A^\mu + \frac{e}{4M} \bar{\psi} (\kappa^S + \kappa^V \tau_0) \sigma_{\mu\nu} \psi F^{\mu\nu}, \quad (8)$$

where $\kappa^S = -0.06$ and $\kappa^V = 1.85$ are the isoscalar and isovector anomalous magnetic moments and $F^{\mu\nu} = \partial^\nu A^\mu - \partial^\mu A^\nu$ we can evaluate the s - and u -channel Born terms. Expressed in the $CGLN$ basis

$$F = iF_1 \vec{\sigma} \cdot \vec{\epsilon} + F_2 \vec{\sigma} \cdot \hat{\vec{q}} \vec{\sigma} \cdot (\hat{\vec{k}} \times \vec{\epsilon}) + iF_3 \vec{\sigma} \cdot \hat{\vec{k}} \hat{\vec{q}} \cdot \vec{\epsilon} + iF_4 \vec{\sigma} \cdot \hat{\vec{q}} \hat{\vec{q}} \cdot \vec{\epsilon} \quad (9)$$

we obtain the following amplitudes for pseudoscalar coupling

$$F_1(PS) = g_\eta C \left[\left(-e_N + \frac{W-M}{2M} \kappa_N \right) D + \frac{(t-m_\eta^2) \kappa_N}{2M(W-M)(u-M^2)} \right], \quad (10a)$$

$$F_2(PS) = g_\eta \frac{C |\vec{q}|}{E_2 + M} \left[\left(e_N + \frac{W+M}{2M} \kappa_N \right) D + \frac{(t-m_\eta^2) \kappa_N}{2M(W+M)(u-M^2)} \right], \quad (10b)$$

$$F_3(PS) = g_\eta C |\vec{q}| \left[2e_N \frac{W-M}{t-m_\eta^2} D - \frac{\kappa_N}{M(u-M^2)} \right], \quad (10c)$$

$$F_4(PS) = g_\eta \frac{C |\vec{q}|^2}{E_2 + M^2} \left[-2e_N \frac{W+M}{t-m_\eta^2} D - \frac{\kappa_N}{M(u-M^2)} \right], \quad (10d)$$

where $t = 2(\vec{k} \cdot \vec{q} - E_\gamma E_\pi) + m_\eta^2$, $u = -2(\vec{k} \cdot \vec{q} + E_\gamma E_2) + M^2$, $E_{1(2)}$ is the nucleon energy in the initial (final) state, and

$$C = -e \frac{W-M}{8\pi W} \sqrt{(E_1+M)(E_2+M)}, \quad D = \frac{1}{W^2-M^2} + \frac{1}{u-M^2}. \quad (11)$$

For pseudovector coupling we get

$$F_1(PV) = F_1(PS) - g_\eta \frac{C \kappa_N}{2M^2}, \quad F_2(PV) = F_2(PS) + g_\eta \frac{C |\vec{q}| \kappa_N}{2M^2(E_2+M)}, \quad (12)$$

and no change for F_3 and F_4 . Note that the difference between pseudoscalar and pseudovector coupling arises only from the anomalous magnetic moment of the nucleon.

In the above equations (where $N = p, n$) the amplitudes are expressed for the protons and neutrons separately with $e_p = 1$, $e_n = 0$ and $\kappa_p = 1.79$, $\kappa_n = -1.91$. Alternatively we can define the isoscalar and isovector amplitudes $F_i^{(0)}$ and $F_i^{(1)}$ via

$$F_i = F_i^{(0)} + F_i^{(1)} \tau_0 \quad (13)$$

Due to the decay of the vector mesons $V(J^\pi; T) = \omega(1^-; 0)$ and $\rho(1^-; 1)$ into $\eta \gamma$ we also have to include the t -channel Born diagrams which we evaluate from the Lagrangians

$$\mathcal{L}_{VNN} = -g_V \bar{\psi} \gamma_\mu \psi V^\mu + \frac{g_T}{4M} \bar{\psi} \sigma_{\mu\nu} \psi V^{\mu\nu}, \quad \mathcal{L}_{V\eta\gamma} = \frac{e\lambda_V}{4m_\eta} \varepsilon_{\mu\nu\lambda\sigma} F^{\mu\nu} V^{\lambda\sigma} \phi_\eta \quad (14)$$

with $V^{\mu\nu} = \partial^\nu V^\mu - \partial^\mu V^\nu$ like the electromagnetic field tensor $F^{\mu\nu}$. This yields to the *CGLN* amplitudes

$$F_1(V) = \frac{\lambda_V C}{m_\eta(t - m_V^2)} \left[-\frac{g_T}{2M} t + \left(\frac{t - m_\eta^2}{2W - 2M} + W - M \right) g_V \right], \quad (15a)$$

$$F_2(V) = \frac{\lambda_V C}{m_\eta(t - m_V^2)} \frac{|\vec{q}|}{E_2 + M} \left[\frac{g_T}{2M} t + \left(\frac{t - m_\eta^2}{2W + 2M} + W + M \right) g_V \right], \quad (15b)$$

$$F_3(V) = \frac{\lambda_V C}{m_\eta(t - m_V^2)} |\vec{q}| \left[\frac{g_T}{2M} (W - M) - g_V \right], \quad (15c)$$

$$F_4(V) = -\frac{\lambda_V C}{m_\eta(t - m_V^2)} \frac{|\vec{q}|^2}{E_2 + M} \left[\frac{g_T}{2M} (W + M) + g_V \right]. \quad (15d)$$

Due to the isospin, the ω contributes only to $F_i^{(0)}$ and ρ only to $F_i^{(1)}$.

In Table II we give the coupling constants and cut-off masses for the background contributions. For the vector mesons we have introduced dipole formfactors $F(\vec{q}^2) = (\Lambda_V^2 - m_V^2)^2 / (\Lambda_V^2 + \vec{q}^2)^2$ at the VNN vertex given by the Bonn potential [16], for the ηNN coupling the formfactors turned out to be insensitive in the energy region of our interest and have been ignored. The main effect would result in a renormalization of the coupling constant. The electromagnetic $V\eta\gamma$ couplings are obtained from the partial decay widths of the vector mesons.

III. THE ηNN COUPLING

In contrast to the πN -interaction, little is known about the ηN -interaction and, consequently, about the ηNN vertex. As it was mentioned before, in the case of pion scattering and pion photoproduction the πNN coupling is preferred to be pseudovector (PV), in accord with current algebra results and chiral symmetry. However, because the eta mass is so much larger than the pion mass - leading to large $SU(3) \times SU(3)$ symmetry breaking - and because of the $\eta - \eta'$ mixing there is no compelling reason to select the PV rather than the PS form for the ηNN vertex.

The uncertainty regarding the structure of the ηNN vertex extends to the magnitude of the coupling constant. This coupling constant $g_{\eta NN}^2/4\pi$ varies between 0 and 7 with the large couplings arising from fits of one boson exchange potentials. Typical values obtained in fits with the Bonn potential [17] can lie anywhere between 3 - 7. However, including the η yields only small effects in fitting the NN phase shifts and, furthermore, provides an insignificant contribution to nuclear binding at normal nuclear densities. From SU(3) flavor symmetry all coupling constants between the meson octet and the baryon octet are determined by one free parameter α , giving

$$\frac{g_{\eta NN}^2}{4\pi} = \frac{1}{3}(3 - 4\alpha)^2 \frac{g_{\pi NN}^2}{4\pi}, \quad (16)$$

resulting in values for the coupling constant between 0.8 and 1.9 for commonly used values of α between 0.6 - 0.65, depending on the F and D strengths chosen as the two types of SU(3) octet meson-baryon couplings. Other determinations of the ηNN coupling employ reactions involving the eta, such as $\pi^- p \rightarrow \eta n$, and range from 0.6 - 1.7 [18]. Smaller values are supported by NN forward dispersion relations [19] with $g_{\eta NN}^2/4\pi + g_{\eta' NN}^2/4\pi \leq 1.0$. There is some rather indirect evidence that also favors a small value for $g_{\eta NN}$. In Ref. [20], Piekarewicz calculated the π - η mixing amplitude in the hadronic model where the mixing was generated by $\bar{N}N$ loops and thus driven by the proton-neutron mass difference. To be in agreement with results from chiral perturbation theory the ηNN coupling had to be constrained to the range 0.32 - 0.53. In a very different approach, Hatsuda [21] evaluated the proton matrix element of the flavor singlet axial current in the large N_C chiral dynamics with an effective Lagrangian that included the $U_A(1)$ anomaly. In this framework, the EMC data on the polarized proton structure function (which have been used to determine the "strangeness content" of the proton) can be related to the $\eta' NN$ and the ηNN coupling constants. Again, his analysis prefers small values for both coupling constants. Nevertheless, from the above discussion it seems clear that the ηNN coupling constant is much smaller compared to the corresponding πNN value of around 14.

Since in our model the resonance sector is well constrained by other related but indepen-

dent reactions we use the (γ, η) data to extract information on the ηNN vertex. In Table III we give the individual contributions to the threshold E_{0+} multipole for protons and neutrons. While the S_{11} contribution is complex, the background contribution from Born terms and vector mesons is real. In the table, the Born PS and PV terms are calculated for a coupling of $g_{\eta NN}^2/4\pi = 1$ and scales proportional to $g_{\eta NN}$. Note that for a coupling of $g_{\eta NN}^2/4\pi = 0.33$ the background vanishes in a pseudoscalar model.

In Fig. 1, we show the sensitivity of the total cross section close to threshold when varying the coupling constant from 0 to 3 for the PS and from 0 to 10 for the PV form. There is a large variation of more than a factor of two at 750 MeV for the PS case while changing the value with PV structure modifies the cross section only by a relatively small amount. The difference is due to the fact that the PV vertex contains momentum dependence which influences mostly the p-wave multipoles. On the other hand, at threshold the total cross section is dominated by the s-wave multipole due to the $S_{11}(1535)$ resonance. A similar effect can be observed in (γ, π^0) at threshold where the PS Born terms overpredict the LET prediction by a large amount while the PV form agrees with the small LET value.

However, using only the total cross section data does not allow to uniquely determine the coupling constant. As shown in Fig. 2, one obtains similar total cross sections for a PS coupling of 0.1 and a PV coupling of 6.0. Data that would fall below this curve could only be explained with a pseudoscalar model. For example, a coupling strength of $g_{\eta NN}^2/4\pi = 0.4$ gives results very close to the pure resonance contribution. Large PS couplings around 1.0 or 1.4 suggested in previous eta photoproduction studies [14] would be consistent only with data considerably below $15 \mu b$ at the maximum. Using only the old data shown in Fig. 2 no definite conclusion can be reached at this point.

Since the total cross section alone cannot unambiguously delineate between the two different coupling modes, we present in Fig. 3 calculations for the differential cross section at the four different photon energies that were used in the Mainz experiment currently under analysis. Note that both computations performed in the PS- and PV-model give roughly the same total cross section. There is a clear distinction in the forward-backward asymmetry of

the angular distribution between the PS- and PV-model. While the older data at 750 MeV (shown in Fig. 3) cannot uniquely distinguish between the two coupling schemes the new preliminary Mainz data with very small error bars indicate a clear preference for a PS-vertex with a small coupling constant. The variation in the angular distributions is again due to the p -wave multipoles. In particular, the M_{1-} multipole changes sign between PS and PV coupling.

IV. MULTIPOLES AND POLARIZATION OBSERVABLES

In order to obtain a detailed understanding of the eta photoproduction process and resonance phenomena associated with it, it would be desirable to perform a multipole analysis along the lines that have been pursued in pion photoproduction for over twenty years. Multipoles offer the possibility to especially study resonance properties in detail since - due to their particular quantum numbers - resonances contribute only to one specific multipole for $J = 1/2$ and to two multipoles (electric and magnetic) for $J \geq 3/2$. In contrast to pion photoproduction, where the $\Delta_{33}(1232)$ resonance dominates for the first 400 MeV above threshold, there are three resonances in eta photoproduction right at threshold - the dominant $S_{11}(1535)$, and the weaker $D_{13}(1520)$ and $P_{11}(1440)$ states. Additional resonances - such as the $P_{11}(1710)$ - are expected to contribute significantly in the region of 200-400 MeV above threshold. Furthermore, the suppression of the Born terms in the (γ, η) process due to the small ηNN coupling constant offers the opportunity to extract valuable resonance information from the (γ, η) multipoles.

In Fig. 4 we present the real and imaginary part of the E_{0+} multipole as a function of the photon energy. Clearly, the S_{11} resonance dominates this multipole and provides the only contribution to the $Im(E_{0+})$ since the Born terms and vector meson contributions - being small and opposite in sign - have no imaginary part. The magnitude of the S_{11} resonance at threshold shows the futility of extracting LET values at threshold. Regarding the ηNN vertex, the E_{0+} is insensitive to the difference between the PS- and PV- coupling by properly

adjusting the coupling constants as was shown in Fig. 2.

This situation changes dramatically for the real part of the M_{1-} multipole shown in Fig. 5. This p -wave multipole which is positive for PS- but negative for PV-coupling is responsible for the variation in the forward-backward asymmetry in the differential cross sections of Fig. 3.

In Fig. 6 we present the recoil polarization P which is proportional to the M_{1-} multipole

$$P = -2\sin(\theta)\frac{q}{k\sigma(\theta)}\text{Im}(E_{0+}^* M_{1-} + \dots) . \quad (17)$$

Thus, the PS-coupling leads to a positive P , supported by a single data point available in this energy region, while PV-coupling leads a negative recoil polarization. Besides the sensitivity to the nature of the ηNN -vertex, Fig. 6 also shows the presence of the $P_{11}(1440)$ state (Roper resonance) in the real part of the M_{1-} multipole. Just as this resonance is not easily identified in other electromagnetic reactions it is not very noticeable in the eta photoproduction process as well.

In Fig. 7 we show the effect of the $D_{13}(1520)$ and the $P_{11}(1440)$ resonances in the differential cross section, $d\sigma/d\Omega$, the three single-polarization observables Σ, T, P and the four double-polarization observables E, F, G, H that require polarization of the beam and the target simultaneously. The energy of 752 MeV is chosen to be in the region where the D_{13} resonance has its maximum contribution. Omitting the D_{13} from our calculations gives dramatic effects in the beam asymmetry Σ as well as in the double-polarization observable G , but T, P, F and H also exhibit a significant sensitivity to the D_{13} state. This is in contrast to the differential cross section that shows very little sensitivity to the D_{13} state. This behavior can be understood in terms of the multipole contributions. Especially the photon asymmetry Σ , which changes from its maximum value around 0.4 at $\theta = 90^\circ$ to almost zero when the $D_{13}(1520)$ is not included, the expansion into leading multipoles yields

$$\Sigma = 3\sin^2(\theta)\frac{q}{k\sigma(\theta)}\text{Re}[E_{0+}^* (E_{2-} + M_{2-}) + \dots] . \quad (18)$$

As in the previous case with the P_{11} , the interference of the D-wave multipoles with the dominant E_{0+} gives such an enhanced sensitivity. While in general all observables of Fig. 7

are different, in this situation with S_{11} dominance (between 700 and 900 MeV) we find similar structures of Σ and G , T and F , and P and H . In future experiments with polarized photon beams and polarized targets comparisons of those pairs of observables can give valuable information on small background multipoles. Furthermore, it can be used to separate real and imaginary parts of the resonance multipoles. Currently, such polarization experiments are already in preparation at LEGS in Brookhaven, GRAAL in Grenoble and CEBAF with an expected start at the end of 1995.

As an example for the $l = 2$ multipoles, Fig. 8 depicts the real part of the E_{2-} multipole. Note that this multipole is again clearly dominated by its resonance contribution since the Born terms and vector mesons almost cancel. Therefore, using total and differential cross section measurements should reveal the strong E_{0+} multipole, while the smaller $l = 1$ and $l = 2$ multipoles could be extracted with the help of polarization experiments.

In the appendix we give the complete structure and definitions of the 16 observables in photoproduction of pseudoscalar mesons as well as the expansion into leading multipoles.

V. ETA PHOTOPRODUCTION ON NUCLEI

Eta photoproduction on nuclei can be developed in a straightforward way by the same method which has been applied very successfully in pion photoproduction [22]. In momentum space the nuclear photoproduction amplitude can be written as

$$F_{\eta\gamma}(\vec{q}, \vec{k}) = V_{\eta\gamma}(\vec{q}, \vec{k}) - \frac{a}{(2\pi)^2} \sum_{i=\pi,\eta} \int \frac{d^3q'}{M_i(q')} \frac{F_{\eta i}(\vec{q}, \vec{q}') V_{i\gamma}(\vec{q}', \vec{k})}{W_\eta(q) - W_i(q') + i\epsilon}, \quad (19)$$

where \vec{k} is the photon, and \vec{q} is the eta or pion momentum. The total energy in the η -nucleus and π -nucleus channels is denoted by $W_i(q) = E_i(q) + E_A(q)$, the reduced mass is given by $M_i(q) = E_i(q)E_A(q)/W_i(q)$ and $a = (A - 1)/A$.

$V_{\eta\gamma}$ is expressed in terms of the free eta-nucleon photoproduction t-matrix

$$V_{\eta\gamma}(\vec{q}, \vec{k}) = -\frac{\sqrt{M_\eta(q)M_\gamma(k)}}{2\pi} \langle \eta(\vec{q}), f | \sum_{j=1}^A \hat{t}_{\gamma N}(j) | i, \gamma(\vec{k}) \rangle, \quad (20)$$

where $|i\rangle$ and $|f\rangle$ denote the nuclear initial and final states, respectively, and j refers to the individual target nucleons.

Using the KMT version of multiple scattering theory [23] the meson scattering amplitude F_{ij} is constructed as a solution of the Lippmann-Schwinger equation

$$F_{ij}(\vec{q}', \vec{q}) = V_{ij}(\vec{q}', \vec{q}) - \frac{a}{(2\pi)^2} \sum_{l=\pi, \eta} \int \frac{d^3 q''}{M_l(q'')} \frac{V_{il}(\vec{q}', \vec{q}'') F_{lj}(\vec{q}'', \vec{q})}{W_j(q) - W_l(q') + i\epsilon}, \quad (21)$$

Here the meson-nuclear interaction is described by the first-order potential $V_{ij} = (V_{\pi\pi}, V_{\eta\pi}, V_{\eta\eta})$ which is related to the corresponding free t_{ij} matrix of meson-nucleon interaction [5].

At present our calculations have been carried out only for the first part of Eq. (19), the plane wave impulse approximation (*PWIA*). At this level, however, we do not perform any approximation treating the full spin degrees of freedom and taking Fermi motion effects of the nucleon into account by performing the integration in momentum space. For the deuteron [24] and $^3\text{He}/^3\text{H}$ [25] we use realistic nuclear wave functions. In the case of ^4He with $J = T = 0$ a phenomenological nuclear formfactor is used which has been extracted from the charge distribution of ^4He .

By studying eta photoproduction on light nuclei with well-known nuclear structure we can learn about details of the elementary production operator which are difficult to see in the elementary reaction or, as for the neutron amplitude, are not experimentally accessible. In the deuteron case only the isoscalar amplitude contributes; in ^3He the two protons saturate to spin 0 and contribute only to a very small part via the non-spin amplitude from the $P_{11}(1440)$ and background terms, while the residual neutron gives rise to a strong E_{0+} amplitude. Finally, in the case of ^4He we can study the coherent amplitude of $t_{\eta\gamma}$ which is the isoscalar non-spin flip part and arises from small magnetic multipoles, e.g. M_{1-} and M_{1+} .

Besides studying the elementary amplitude, eta photoproduction offers the possibility to learn more about the ηN - and ηA -interaction. Recently it has been suggested by Wilkin [26] that the very large production cross sections found in the $pd \rightarrow \eta^3\text{He}$ reaction near threshold

could be due to an ηN scattering length that is much larger than the value extracted from a coupled-channels analysis of the $\pi^- p \rightarrow \eta n$ and $\pi N \rightarrow \pi N$ data [27]. Based on the K-matrix formalism, Wilkin was able to reproduce the strong threshold energy dependence of $\pi d \rightarrow \eta^3\text{He}$ by assuming an ηN scattering length with a real part more than twice as large as in previous analyses. Should this conclusion turn out to be true, it would have dramatic implications for the η -nucleus interaction. Most importantly, a larger value for the scattering length might indicate a larger probability for the presence of a "bound" η -state for lighter nuclei such as ^3He than had been expected. In fact, this seems to be required to explain the cusp-like structure seen for near-threshold production in the $\pi d \rightarrow ^3\text{He} X$ reaction for a missing mass close to the η -mass. An uncertainty in Wilkin's analysis arises from a cross section measurement of the $pd \rightarrow \eta^3\text{He}$ process very close to threshold (~ 200 keV) that seems to contradict his description of the energy dependence. However, if this one point - which suffers from large systematic uncertainties - is ignored, impressive agreement with the experiment is achieved. A new experiment at Saclay has already been approved to explore the energy dependence in the region very close to threshold [28]. This effect may also remove the discrepancy between experiment and theory for the reaction $^3\text{He}(\pi^-, \eta)^3\text{H}$ [5]. In a future study we will address the significance of this large $\eta^3\text{He}$ -interaction in the $^3\text{He}(\gamma, \eta)^3\text{He}$ cross section calculated in a full DWIA framework.

In Fig. 9 we show the differential cross sections for all light nuclei up to ^4He . Whereas the angular distribution is rather flat for nucleons, as shown before, it appears more and more peaked in forward direction for $A > 1$. This reflects the signature of the nuclear formfactors as the momentum transfer in η photoproduction is rather large, $Q^2 = 7.8 \text{ fm}^{-2}$ at threshold. The biggest cross section can be expected for the trinucleon; it is proportional to the free nucleon cross section multiplied by the square of the trinucleon formfactor. The triton cross section is about a factor of two larger than that for ^3He since the spin-flip amplitude on the proton - which is the largest among the possible spin-flip and non-spin-flip amplitudes on the proton and the neutron - cannot contribute due to the Pauli principle. Therefore, the single proton in ^3H provides more strength than the two protons (coupled mostly to spin 0)

in ^3He .

Around 90° the cross section on the deuteron gains over the trinucleon, even though it is in large disagreement with the few available data. This may be due to the rather small isoscalar amplitude predicted by our model, $E_{0+}^{(0)}/E_{0+}^{(p)} = 0.15$. In the naive quark model the electromagnetic excitation of the S_{11} is almost entirely isovector. However, Rosenthal, Forest and Gonzales [29] have shown that a color-hyperfine interaction, responsible also for the E2/M1 ratio of the Δ excitation, could enhance the isoscalar S_{11} excitation considerably.

Only an unrealistically large isoscalar amplitude of $E_{0+}^{(0)}/E_{0+}^{(p)} = 0.8$ can explain the deuteron data in *PWIA*. As it has been shown in Ref. [30] if the final state interaction and the coupled $\pi\eta$ -channels are taken into account then the discrepancy can be explained with $E_{0+}^{(0)}/E_{0+}^{(p)} = 0.6$. It also remains to be seen if the isoscalar amplitude is really as small as all present models predict. Finally, the coherent cross section for ^4He vanishes for $\theta = 0$ and reaches roughly the 10nb level in a small angular region. For most angles it falls below 1nb . It is an experimental challenge to measure this reaction which provides a clean observation of the background multipoles and of the Roper resonance, as the dominant S_{11} resonance is suppressed by spin and isospin.

VI. SUMMARY AND CONCLUSIONS

We have presented a model for eta photoproduction on the nucleon that includes nucleon Born terms and t-channel vector meson exchanges in addition to the nucleon resonances $S_{11}(1535)$, $P_{11}(1440)$ and $D_{13}(1520)$. The resonance sector is fixed by using data from other hadronic and electromagnetic reactions such as pion scattering and photoproduction and pion induced eta production. Vector meson couplings are determined from their radiative decay widths and the NN -interaction. This allows using the new experimental data from Bonn and Mainz to extract information on the ηNN coupling. While the total (γ, η) cross section on the proton can be well reproduced by either a small coupling constant with PS-coupling or a large value with PV-form, the angular distribution singles out the small

constant.

While the resonance parameters of the S_{11} can be extracted from high-precision total cross section data, the smaller resonances are hidden both in the differential and total cross sections. Here, polarization observables will provide a powerful tool to constrain these small resonance couplings, as demonstrated for the case of the $D_{13}(1520)$. Such experiments have been proposed and will soon be performed at the LEGS facility at Brookhaven National Laboratory and at the new facility GRAAL at Grenoble.

Finally, we have applied our operator to eta photoproduction on light nuclei d, ^3He , ^3H and ^4He . Experiments on these nuclei are necessary to obtain complete information on the isospin structure of the (γ, η) amplitude. At forward angles the cross sections for (γ, η) on the trinucleon are reasonably large and should be measurable. Here especially the threshold region would be interesting where one could learn more about the eta-nucleus interaction.

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APPENDIX A: POLARIZATION OBSERVABLES

Following the notation of Barker et al. [31] we can write the differential cross section

a) for beam and target polarization

$$\begin{aligned} \frac{d\sigma}{d\Omega} = & \sigma_0 \{ 1 - P_T \Sigma \cos 2\varphi \\ & + P_x (-P_T H \sin 2\varphi + P_\odot F) - P_y (-T + P_T P \cos 2\varphi) \\ & - P_z (-P_T G \sin 2\varphi + P_\odot E) \} , \end{aligned} \tag{A1}$$

b) for beam and recoil polarization

$$\begin{aligned} \frac{d\sigma}{d\Omega} = & \sigma_0 \{ 1 - P_T \Sigma \cos 2\varphi \\ & + P_{x'} (-P_T O_{x'} \sin 2\varphi - P_\odot C_{x'}) - P_{y'} (-P + P_T T \cos 2\varphi) \\ & - P_{z'} (P_T O_{z'} \sin 2\varphi + P_\odot C_{z'}) \} , \end{aligned} \quad (\text{A2})$$

c) for target and recoil polarization

$$\begin{aligned} \frac{d\sigma}{d\Omega} = & \sigma_0 \{ 1 + P_{y'} P + P_x (P_{x'} T_{x'} + P_{z'} T_{z'}) \\ & + P_y (T + P_{y'} \Sigma) - P_z (P_{x'} L_{x'} - P_{z'} L_{z'}) \} , \end{aligned} \quad (\text{A3})$$

where P_T and P_\odot denote the degree of linear and right-handed circular photon polarization. (P_x, P_y, P_z) is the polarization of the target in the right-handed frame $\{x, y, z\}$, with \hat{z} along the photon axis and $\hat{y} = \hat{k}_\gamma \times \hat{q}_\pi / \sin \theta$. The spin of the recoil nucleon is analyzed as $(P_{x'}, P_{y'}, P_{z'})$ in the frame $\{x', y', z'\}$ with \hat{z}' along the meson axis and $\hat{y}' = \hat{y}$. The angle θ of the meson as well as all other quantities are measured in the *cm* frame. The azimuthal angle φ of the vector of linear photon polarization is measured counter-clockwise from the scattering plane, e.g. $\varphi = \pi/2$ for a photon polarization ϵ_\perp along the \hat{y} -axis. The unpolarized cross section σ_0 will be expressed in terms of the transverse response function R_T ,

$$\sigma_0 \equiv \left. \frac{d\sigma}{d\Omega} \right|_{\text{unpolarized}} = \frac{q}{k} R_T . \quad (\text{A4})$$

APPENDIX B: CGLN AMPLITUDES

The 16 polarization observables for pseudoscalar meson photoproduction can be expressed in terms of the four complex CGLN amplitudes of Eq. (9):

$$\begin{aligned} R_T = & \text{Re} \{ |F_1|^2 + |F_2|^2 - 2 \cos \theta F_1^* F_2 + \\ & + \frac{\sin^2 \theta}{2} (|F_3|^2 + |F_4|^2 + 2F_2^* F_3 + 2F_1^* F_4 + 2 \cos \theta F_3^* F_4) \} , \end{aligned}$$

$$\begin{aligned}
R_T \Sigma &= -\frac{\sin^2 \theta}{2} \text{Re}\{|F_3|^2 + |F_4|^2 + 2(F_2^* F_3 + F_1^* F_4 + \cos \theta F_3^* F_4)\}, \\
R_T T &= \sin \theta \text{Im}\{F_1^* F_3 - F_2^* F_4 + \cos \theta (F_1^* F_4 - F_2^* F_3) - \sin^2 \theta F_3^* F_4\}, \\
R_T P &= \sin \theta \text{Im}\{F_2^* F_4 - 2F_1^* F_2 - F_1^* F_3 + \cos \theta (F_2^* F_3 - F_1^* F_4) + \sin^2 \theta F_3^* F_4\}, \\
R_T G &= \sin^2 \theta \text{Im}\{F_2^* F_3 + F_1^* F_4\}, \\
R_T H &= \sin \theta \text{Im}\{2F_1^* F_2 + F_1^* F_3 - F_2^* F_4 - \cos \theta (F_2^* F_3 - F_1^* F_4)\}, \\
R_T E &= \text{Re}\{|F_1|^2 + |F_2|^2 - 2 \cos \theta F_1^* F_2 + \sin^2 \theta (F_2^* F_3 + F_1^* F_4)\}, \\
R_T F &= \sin \theta \text{Re}\{F_1^* F_3 - F_2^* F_4 + \cos \theta (F_1^* F_4 - F_2^* F_3)\}, \\
R_T O_{x'} &= -\sin \theta \text{Im}\{F_1^* F_4 - F_2^* F_3 + \cos \theta (F_1^* F_3 - F_2^* F_4)\}, \\
R_T O_{z'} &= -\sin^2 \theta \text{Im}\{F_1^* F_3 + F_2^* F_4\}, \\
R_T C_{x'} &= \sin \theta \text{Re}\{|F_1|^2 - |F_2|^2 + F_1^* F_4 - F_2^* F_3 + \cos \theta (F_1^* F_3 - F_2^* F_4)\}, \\
R_T C_{z'} &= \text{Re}\{2F_1^* F_2 + \sin^2 \theta (F_1^* F_3 + F_2^* F_4) - \cos \theta (|F_1|^2 + |F_2|^2)\}, \\
R_T T_{x'} &= -\sin^2 \theta \text{Re}\{\frac{\cos \theta}{2}(|F_3|^2 + |F_4|^2) + F_1^* F_3 + F_2^* F_4 + F_3^* F_4\}, \\
R_T T_{z'} &= \sin \theta \text{Re}\{\frac{\sin^2 \theta}{2}(|F_4|^2 - |F_3|^2) + F_1^* F_4 - F_2^* F_3 + \cos \theta (F_1^* F_3 - F_2^* F_4)\}, \\
R_T L_{x'} &= \sin \theta \text{Re}\{\frac{\sin^2 \theta}{2}(|F_3|^2 - |F_4|^2) - |F_1|^2 + |F_2|^2 + \\
&\quad + F_2^* F_3 - F_1^* F_4 - \cos \theta (F_1^* F_3 - F_2^* F_4)\}, \\
R_T L_{z'} &= \text{Re}\{2F_1^* F_2 + \sin^2 \theta (F_1^* F_3 + F_2^* F_4 + F_3^* F_4) - \\
&\quad - \cos \theta (|F_1|^2 + |F_2|^2) + \frac{\sin^2 \theta}{2} \cos \theta (|F_3|^2 + |F_4|^2)\}.
\end{aligned}$$

APPENDIX C: EXPANSION IN LEADING MULTIPOLES

Complete expansions of the polarization observables into multipoles up to $l = 1$ can be found in Refs. [32,33].¹ As this is sufficient for most cases of pion photoproduction, in eta photoproduction also $l = 2$ multipoles play an important role even at threshold. However, a

¹Note that Ref. [33] uses a different sign for the observables $E, H, O_{x'}, O_{z'}, C_{x'}, C_{z'}$ and $L_{x'}$.

general expansion up to $l = 2$ is very extensive and difficult to survey. Therefore, we give only those leading multipoles that are excited by nucleon resonances, the E_{0+} by the $S_{11}(1535)$, the M_{1-} by the $P_{11}(1440)$ and the E_{2-}, M_{2-} excited by the $D_{13}(1520)$. Furthermore, since the S_{11} dominates so strongly, we restrict ourselves only to those contributions proportional to E_{0+} .

$$\begin{aligned}
R_T &= |E_{0+}|^2 - \text{Re}\{E_{0+}^*[2\cos\theta M_{1-} - (3\cos^2\theta - 1)(E_{2-} - 3M_{2-})]\} , \\
R_T\Sigma &= 3\sin^2\theta \text{Re}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
R_T T &= -3\sin\theta \cos\theta \text{Im}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
R_T P &= -\sin\theta \text{Im}\{E_{0+}^*[2M_{1-} - 3\cos\theta(E_{2-} - 3M_{2-})]\} , \\
R_T G &= -3\sin^2\theta \text{Im}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
R_T H &= \sin\theta \text{Im}\{E_{0+}^*[2M_{1-} - 3\cos\theta(E_{2-} - 3M_{2-})]\} , \\
R_T E &= |E_{0+}|^2 - \text{Re}\{E_{0+}^*[2\cos\theta M_{1-} - (3\cos^2\theta - 1)(E_{2-} - 3M_{2-})]\} , \\
R_T F &= -3\sin\theta \cos\theta \text{Re}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
R_T O_{x'} &= 3\sin\theta \text{Im}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
O_{z'} &= 0 , \\
R_T C_{x'} &= \sin\theta [|E_{0+}|^2 - \text{Re}\{E_{0+}^*(E_{2-} - 3M_{2-})\}] , \\
R_T C_{z'} &= -\cos\theta |E_{0+}|^2 + 2\text{Re}\{E_{0+}^*[M_{1-} - \cos\theta(E_{2-} - 3M_{2-})]\} , \\
T_{x'} &= 0 , \\
R_T T_{z'} &= -3\sin\theta \text{Re}\{E_{0+}^*(E_{2-} + M_{2-})\} , \\
R_T L_{x'} &= -\sin\theta [|E_{0+}|^2 - \text{Re}\{E_{0+}^*(E_{2-} - 3M_{2-})\}] , \\
R_T L_{z'} &= -\cos\theta |E_{0+}|^2 + 2\text{Re}\{E_{0+}^*[M_{1-} - \cos\theta(E_{2-} - 3M_{2-})]\} .
\end{aligned}$$

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FIGURES

FIG. 1. Total cross section for the process (γ, η) on the proton calculated with PS and PV Born terms. The full curve contains no Born terms, while the dashed lines are (from the top down) obtained with $g_{\eta NN}^2/4\pi=0.1, 0.5, 1.0$, and 3.0 for PS-coupling, and $g_{\eta NN}^2/4\pi=1.0, 3.0, 6.0$ and 10.0 for PV-coupling, respectively.

FIG. 2. Total cross section for the process (γ, η) on the proton. The full (dashed) curve is obtained with resonances, vector mesons and PS (PV) Born terms with $g_{\eta NN}^2/4\pi = 0.1$ (6.0). The dotted curve shows the nucleon resonances only while the dashed-dotted line also includes vector mesons but not Born terms. The experimental data are from Ref. [7] (\bullet) and Ref. [34] (\circ).

FIG. 3. Differential cross section for the process (γ, η) on the proton. The curves are as in Fig. 2. The experimental data are from Ref. [7] (\bullet) and Ref. [6] (\circ).

FIG. 4. Real and imaginary parts of the E_{0+} multipole for (γ, η) on the proton. The curves are as in Fig. 2. The resonance contribution for this multipole comes solely from the $S_{11}(1535)$.

FIG. 5. Real part of the M_{1-} multipole for (γ, η) on the proton. The curves are as in Fig. 2. The resonance contribution for this multipole comes solely from the Roper.

FIG. 6. Recoil polarization for (γ, η) on the proton at a photon lab energy of 830 MeV. The curves are as in Fig. 2. The experimental data point is from Ref. [35].

FIG. 7. Influence of the $P_{11}(1440)$ and $D_{13}(1520)$ resonances on the differential cross section $d\sigma/d\Omega$, the single-polarization observables Σ , T and P and the double-polarization observables E, F, G, H for polarization of beam and target at a photon lab energy of 752 MeV. The full lines show the complete calculation with resonances, vector mesons and PS Born terms with $g_{\eta NN}^2/4\pi = 0.4$. The dashed and dotted lines are obtained when the D_{13} or the P_{11} resonances are omitted, respectively.

FIG. 8. Real part of the E_{2-} multipole for (γ, η) on the proton. The curves are as in Fig. 2.

FIG. 9. Differential cross section for eta photoproduction on p, n, d, ^3He , ^3H and ^4He . The experimental data on the proton are from Ref. [7] (\bullet) and Ref. [6] (\circ), the data point on the deuteron is from Ref. [36] (\triangle).

TABLES

TABLE I. Parameters for the πN and ηN self-energies in MeV

		a	b_1	b_2	c_1	c_2	d_1	d_2	d_3
S_{11}	πN	17	0	0	0	0	-129.5	80	-5
	ηN	-27	17.7	-1.23	22.9	-5.17	-38.1	18.3	0
P_{11}	πN	-150	0	0	0	0	55.1	-96.2	6.6
D_{13}	πN	-26	0	0	0	0	23	-32.1	2.7

TABLE II. Coupling constants and cut-off masses for the background vector meson exchange contributions.

V	$g_V^2/4\pi$	g_T/g_V	$\Lambda_V(MeV)$	λ_V
ω	23	0	1400	0.192
ρ	0.5	6.1	1800	0.89

TABLE III. Contributions to the threshold amplitudes of E_{0+} in units of $10^{-3}/m_\pi$. In the lab frame, the threshold photon energies are $707.16 MeV$ on protons and $706.94 MeV$ on neutrons for an eta mass of $547.45 MeV$. The Born terms have been calculated with $g_{\eta NN}^2/4\pi=1.0$

target	$S_{11}(1535)$	ω	ρ	Born $PS(1.0)$	Born $PV(1.0)$
proton	12.91 + 5.97 i	0.35	2.63	-5.20	-0.88
neutron	-7.12 - 4.86 i	0.35	-2.63	3.55	-1.04

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